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Brane Mediated Supersymmetry Breaking

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Abstract

We propose a mechanism for mediating supersymmetry breaking in Type I string constructions. The basic set-up consists of a system of three D-branes: two parallel D-branes, a matter D-brane and a source D-brane, with supersymmetry breaking communicated via a third D-brane, the mediating D-brane, which intersects both of the parallel D-branes. We discuss an example in which the first and second family matter fields correspond to open strings living on the intersection of the matter D-brane and mediating D-brane, while the gauge fields, Higgs doublets and third family matter fields correspond to open strings living on the mediating D-brane. As in gaugino mediated models, the gauginos and Higgs doublets receive direct soft masses from the source brane, and flavour-changing neutral currents are naturally suppressed since the first and second family squarks and sleptons receive suppressed soft masses. However, unlike the gaugino mediated model, the third family squarks and sleptons receive unsuppressed soft masses, resulting in a very distinctive spectrum with heavier stops, sbottoms and staus.

1 Introduction

The process of SUSY breaking continues to be an active area of research. Over the years there have been various mechanisms proposed, including gauge [1] and anomaly mediation [2]. An alternative mechanism has been been put forward [3, 4] called *gaugino* mediated supersymmetry (SUSY) breaking which has the attractive property of solving the flavour problem, since scalars masses effectively vanish at the GUT scale and are generated through radiative corrections for which a GIM-like mechanism prevents flavour-changing neutral current (FCNC) problems. This is rather like the no-scale supergravity mechanism [5], but is implemented within a Horava-Witten [6] type set-up [1] consisting of two parallel but spatially separated D3-branes with SUSY broken on one brane, with the SUSY matter fields living on the other brane and the gauge sector living in the bulk and communicating the SUSY breaking from one brane to the other. The Higgs doublets may also be in the bulk providing a solution to the μ problem via the Giudice-Masiero mechanism [7]. The advantage of this set-up is that the contact terms arising from integrating out states with mass M are suppressed by a Yukawa factor e^{-Mr} if $M \geq r$, and so a modest separation between the two branes can lead to negligible direct communication between the SUSY breaking brane and the matter brane. This is the starting point of both the anomaly mediated and the gaugino mediated models, and underpins the solution to the FCNC problem in both cases.

In this paper we shall propose a mechanism for mediating SUSY breaking in Type I string models based on open strings starting and ending on D-branes. Type I string theories can provide an attractive setting for ideas such as gaugino mediated SUSY breaking ($\tilde{g}MSB$), and we shall explore this possibility in this paper. In place of the Horava-Witten set-up we shall consider a Type I toy model consisting of two parallel D-branes with a third D-brane intersecting with both of the parallel D-branes. Instead of having the gauge fields in the bulk we shall put the gauge fields onto the third mediating D-brane, which allows SUSY breaking to be communicated between the SUSY breaking brane and the matter brane. Thus the role of the bulk is played by

[1] Note that in the Horava-Witten model gauge fields do not live in the bulk.

the third mediating D-brane, and it is the gauge fields which live on this brane that communicates the SUSY breaking. However in Type I models it is natural for a matter family to also live on the mediating D-brane, and this provides a characteristic signature of the brane mediated SUSY breaking mechanism.

To illustrate these ideas we consider a toy model inspired by the work of Shiu and Tye [8] using intersecting D5-branes, where the intersection regions are effectively parallel D3-branes within a higher-dimensional spacetime. In this model two chiral families occur in the 4d intersection region at the origin fixed point ($5_1 5_2$ sector), with a third family on the $D5_2$-brane ($5_2 5_2$ sector). However our model differs from Shiu-Tye since we include a further $D5_1'$-brane which intersects with the $D5_2$-brane at a point located away from the origin fixed point, and suppose that SUSY gets broken on that brane and is communicated via the states on the $D5_2$-brane which intersect with both $D5_1$-branes at the two fixed points – brane mediated SUSY breaking (BMSB). In this example gauge fields, Higgs fields and the third family all live in the mediating D-brane which plays the role of the bulk in the original scenario. This separation of the third family[2] provides an explanation for the large mass of the third family of quarks and leptons, without perturbing the solution to the flavour problem since the first and second families remain almost degenerate.

The layout of the remainder of the paper is as follows. In section 2 we review $\tilde{g}MSB$, and in section 3 we introduce a Type I string-inspired toy model motivated by Shiu and Tye. Section 4 is the main section of the paper in which we present our toy model that illustrates the BMSB mechanism, and explore its theoretical and experimental consequences. Section 5 concludes the paper.

2 Gaugino Mediated Supersymmetry Breaking

In this section we review the $\tilde{g}MSB$ mechanism in Refs. [3, 4]. This toy model involves D3-branes embedded in a higher-dimensional space. Two parallel D3-branes are spatially separated along (at least) one extra dimension as shown in Fig. 1. Standard Model quark and lepton fields

[2]Remember that the first two families are localised within an effective 4d overlapping region, while the third family feels two extra dimensions

are localised on the matter brane as open strings, while the gauge and (possibly) Higgs fields propagate in the bulk[3]. Supersymmetry is broken on the displaced source D3-brane. SUSY breaking is communicated to the bulk fields by direct higher-dimensional interactions[4], and mediated to the quark/lepton fields by Standard Model loops[5].

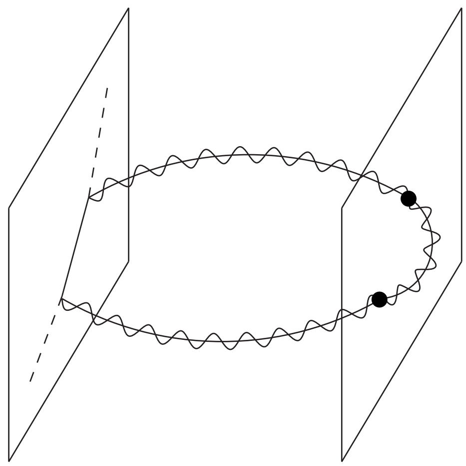

Figure 1: An extra dimensional loop diagram that contributes to SUSY breaking scalar masses. It is similar to a self-energy diagram, but with the virtual gaugino not confined to either 4-dimensional brane. This Figure is taken from Ref. [3].

The full D-dimensional lagrangian is split into two distinct pieces - a bulk term involving only bulk fields and terms localised on either D3-brane that allow direct bulk-brane field coupling.

$$\mathcal{L}_D = \mathcal{L}_{bulk}\left(\Phi_{bulk}\left(x,y\right)\right) + \sum_j \delta^{D-4}\left(y-y_j\right)\mathcal{L}_j\left(\Phi_{bulk}\left(x,y_j\right),\phi_j\left(x\right)\right) \tag{1}$$

where j runs over the branes, x are coordinates for the 4 non-compact dimensions, y are coordinates for the $D-4$ compact spatial dimensions, Φ_{bulk} is a bulk field, and ϕ_j is a field localised on the j^{th} brane.

A Naive Dimensional Analysis (NDA) allows the 5d (or higher) effective theory to be *matched*

[3]Thus feeling all 5-dimensions.

[4]Higher-dimensional operators are assumed to arise from the underlying string theory, although this is not clear at present.

[5]Gauginos in the bulk couple directly to chiral fermions on the matter brane. They also couple to the hidden sector directly through mass-insertions on the source brane.

on to the observed 4d theory at the compactification scale. The 4 and D-dimensional gauge couplings can be related by the size of the compact dimensions:

$$g_4^2 = \frac{g_D^2}{V_{D-4}} \tag{2}$$

The D-dimensional gauge coupling g_D must be smaller than its strong-coupling limit, otherwise perturbative results become meaningless[6]

$$g_D^2 \sim \frac{\epsilon l_D}{M^{D-4}} \tag{3}$$

where l_D is a geometrical loop factor for D dimensions, $l_D = 2^D \pi^{D/2} \Gamma(D/2)$, M is the fundamental scale in the theory which acts as a regulating cutoff, and ϵ suppresses the coupling strength. $\epsilon \sim 1$ corresponds to the strong coupling limit. This places a constraint, along with FCNC suppression, that restricts the maximum size of the extra dimensions. (See [3, 4] for details.)

Following the work of Randall and Sundrum on spatially-separated D3-branes in extra dimensions [2], contact terms between fields on opposite branes are exponentially suppressed by an amount e^{-ML}, where L is the separation between D3-branes along the extra dimension(s).

Eq. (4) is an example of an exponentially suppressed 4-point operator involving fields from the matter and source branes that generates scalar masses:

$$\triangle \mathcal{L}_{brane} \sim \frac{e^{-ML}}{M^2} \int d^4\theta \left(\hat{\phi}_S^\dagger \hat{\phi}_S \right) \left(\phi_M^\dagger \phi_M \right) \tag{4}$$

(where ϕ_S, ϕ_M are source and matter fields respectively)

Compare the suppressed contact terms with the operators giving rise to gaugino masses and Higgs SUSY breaking parameters, from Higgs superfields h_u, h_d and gauge field strengths W_α living in the bulk[7].

$$\begin{aligned}\triangle \mathcal{L}_{brane} \sim \frac{l_D}{l_4} \left(\int d^2\theta \frac{1}{M^{D-3}} \hat{\phi}_S W^\alpha W_\alpha + h.c. \right) + \frac{l_D}{l_4} \int d^4\theta \left\{ \frac{1}{M^{D-3}} \left(\hat{\phi}_S^\dagger h_u h_d + h.c. \right) \right. \\ \left. + \frac{1}{M^{D-2}} \hat{\phi}_S^\dagger \hat{\phi}_S \left[h_u^\dagger h_u + h_d^\dagger h_d + (h_u h_d + h.c.) \right] \right\}\end{aligned} \tag{5}$$

[6]Extra dimensions (and Kaluza-Klein excitations) change the energy-dependence of couplings to power law running above the compactification scale. This allows for unification at lower scales, see [9] for a review.

[7]The scale factors M arise from the requirement of canonical normalization.

This leads to soft terms when we match to the D-dimensional theory[8] and using eqs. (2,3) with $g_4 \approx 1$:

$$m_\lambda, \mu \sim \frac{\hat{F}_S}{M}\frac{l_D/l_4}{M^{D-4}V_{D-4}} \sim \frac{1}{\epsilon l_4}\frac{\hat{F}_S}{M}, \qquad B\mu, m_{h_u}^2, m_{h_d}^2 \sim \frac{\hat{F}_S^2}{M^2}\frac{l_D/l_4}{M^{D-4}V_{D-4}} \sim \frac{1}{\epsilon l_4}\frac{\hat{F}_S^2}{M^2} \tag{6}$$

Both papers discuss methods of generating the μ-term[9]. Ref. [3] suggested the inclusion of an additional gauge singlet on the matter brane (NMSSM) with an extra superpotential term W $\sim \lambda N h_u h_d$. An effective μ-term is produced if N acquires a non-zero vacuum expectation value (vev). Another possibility [4] is to produce the μ-term on the source brane through the Giudice-Masiero mechanism [7] (as above) $\mathcal{L} \sim \int d^4\theta \lambda_\mu \hat{\phi}_S^\dagger h_u h_d$.

3 Type I String-Inspired Model

Now we turn to Type I string constructions and introduce a toy model motivated by the work of Shiu and Tye [8]. The string scale m_s is usually considered to be of the order 10^{16} GeV, but recently the gauge unification scale was suggested to be as low as 1 TeV, which could allow the string scale at a comparable value. Shiu and Tye [8] discuss the phenomenological possibilities within Type I string theory and overlapping D5-branes. They use the duality between the compactification of 10-dimensional Type IIB string theory on an orientifold, with Type I theory on an orbifold to recover a 4-dimensional $\mathcal{N} = 1$ supersymmetric chiral string model with Pati-Salam-like gauge symmetry.

Tadpole cancellations and a non-zero background NS-NS B-field constrain the number and type of D-branes allowed within the model to D5 and D9-branes only [10]. In a particular scenario they consider only one type of D5 brane (5_3) together with the D9 brane, and after T-dualizing they arrive at a scenario with two intersecting branes, namely 5_1 and 5_2 branes which intersect at the origin fixed point. A gauge group $U(4) \otimes U(2) \otimes U(2)'$ exists on each

[8] Notice that the $B\mu$ term and Higgs mass-squared terms are enhanced by a volume factor relative to the m_λ, μ terms.

[9] Ref. [3] has the Higgs fields localised on the matter D3-brane, while [4] has the Higgs fields living in the bulk.

brane, and they discuss three scenarios where the Standard Model gauge group originates from different brane combinations. Their third scenario is of particular interest since it leads to three chiral families - two families on the $5_1 5_2$ overlap and a third family on the $D5_2$-brane as shown in Fig. 2.

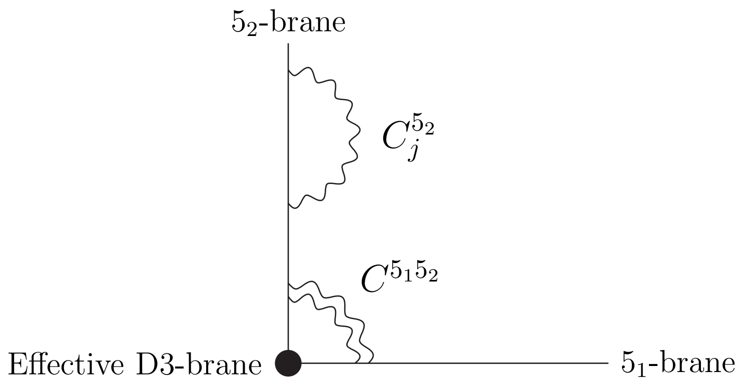

Figure 2: The matter fields and Higgs doublets resulting from Shiu and Tye's third scenario with intersecting D5-branes, where C_i^p is an open-string state (matter field) starting and ending on the p^{th} brane. C^{pq} is an open-string state starting on the p^{th} brane and ending on the q^{th} brane.

We can express the allowed superpotential [11] in terms of the possible states from the two types of D5-branes present in this model:

$$W = C_1^{5_2} C_2^{5_2} C_3^{5_2} + C_3^{5_2} C^{5_1 5_2} C^{5_1 5_2} \tag{7}$$

We now proceed to introduce a toy model based on the above construction. [10] In order to allow the third family Yukawa couplings $\left(\bar{F}_3 F_3 h\right)$ consistent with the string selection-rules in eq. (7), we shall assign the Higgs $h_u, h_d \equiv C_1^{5_2}$ or $C_2^{5_2}$. This leads to the four possible allocations of 5_2 states in Table 1.

Notice that there are no free indices on the intersection states $Q_i, L_i, U_i^C, D_i^C, E_i^C, N_i^C \equiv C^{5_1 5_2}$ (i=1,2), which means that we cannot distinguish between the first two families.

In our Type I string-inspired model, we shall assign the gauge groups and matter fields as in Table 2. We ignore the custodial $SU(4)_{5_1} \otimes U(1)^6$ symmetry. The states $\phi, \phi', H^{\alpha b}, \bar{H}_{\alpha b}$ are used to break the gauge group down to the Standard Model as discussed in Appendix A.1.

[10] For other examples of toy models based on this construction see [12] and references therein.

5_2 states	A	B	C	D
$h \sim h_u, h_d$	$C_1^{5_2}$	$C_1^{5_2}$	$C_2^{5_2}$	$C_2^{5_2}$
$F_3 \sim Q_3, L_3$	$C_2^{5_2}$	$C_3^{5_2}$	$C_3^{5_2}$	$C_1^{5_2}$
$\bar{F}_3 \sim U_3^C\ ,\ D_3^C\ ,\ E_3^C\ ,\ N_3^C$	$C_3^{5_2}$	$C_2^{5_2}$	$C_1^{5_2}$	$C_3^{5_2}$

Table 1: Allocation of 5_2 states that lead to third family-only Yukawa couplings at lowest order. We use the lower index to distinguish between doublets, singlets and Higgs fields.

States	Sector	$SU(4)_{5_2}$	$SU(2)_{5_{2R}}$	$SU(2)_{5_{2L}}$	$SU(2)_{5_{1R}}$	$SU(2)_{5_{1L}}$
$F_i \sim Q_i, L_i$	$5_1 5_2$	4	1	1	1	2
$\bar{F}_i \sim U_i^C, D_i^C, E_i^C, N_i^C,$	$5_1 5_2$	$\bar{4}$	1	1	2	1
$F_3 \sim Q_3, L_3$	5_2	4	1	2	1	1
$\bar{F}_3 \sim U_3^C, D_3^C, E_3^C, N_3^C,$	5_2	$\bar{4}$	2	1	1	1
ϕ	$5_1 5_2$	1	1	2	1	2
ϕ'	$5_1 5_2$	1	2	1	2	1
$H^{\alpha b}$	5_2	4	2	1	1	1
$\bar{H}_{\alpha b}$	5_2	$\bar{4}$	2	1	1	1
$h \sim h_u, h_d$	5_2	1	2	2	1	1

Table 2: $SU(4)_{5_2} \otimes SU(2)_{5_{2R}} \otimes SU(2)_{5_{2L}} \otimes SU(2)_{5_{1R}} \otimes SU(2)_{5_{1L}}$ quantum numbers for left and right-handed chiral fermion states and symmetry breaking Higgs fields.

Gauge invariance with respect to the initial gauge group $SU(4)_{5_2} \otimes SU(2)_{5_{2R}} \otimes SU(2)_{5_{2L}} \otimes SU(2)_{5_{1R}} \otimes SU(2)_{5_{1L}}$ provides the mechanism to forbid both first and second family Yukawa couplings $\left(\bar{F}_i F_j h\right)$ and R-parity violating operators without any other assumptions[11].

Note that the μ-term is forbidden by string selection rules which also forbid a superpotential term involving a matter brane singlet[12] where W $\sim \lambda N h_u h_d$. The Giudice-Masiero mechanism offers the best opportunity of producing a μ-term from the soft potential as discussed later.

[11] Note that the third family right-handed neutrinos and sneutrinos receive large Majorana masses from the operators $\bar{F}_3 \bar{F}_3 H H$ resulting in a see-saw mechanism. This is discussed in Ref. [13], along with a discussion of higher-dimensional operators suitable for first and second family fermion masses.

[12] A non-renormalisable higher-dimensional 4-point superpotential term may be generated by two additional gauge singlet fields, eg. $W \sim N_1 N_2 h_u h_d$. This can become the 3-point term when one of the singlet fields acquire a vev.

4 Brane Mediated Supersymmetry Breaking

We now augment the model in section 3, including the states in Table 2, by including an additional $D5'_1$-brane located at an orbifold fixed point away from the origin as shown in Fig. 3. The idea of including the extra $5'_1$-brane is that SUSY is broken on this brane and communicated by the MSSM states that live on the 5_2-brane which intersects it. Thus, the gauge fields on the 5_2-brane play the role of the gauge fields in the bulk in Fig. 1. Note that there are many mass scales in this model as discussed in Appendix B.

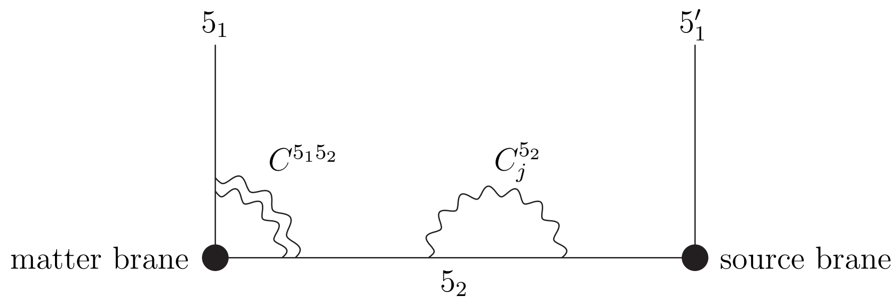

Figure 3: A brane-construction using overlapping D5-branes, with effective D3-branes at the intersection points spatially separated along the $D5_2$-brane. The first two chiral families ($C^{5_1 5_2}$) live on the first intersection region. The third family and Higgs doublets ($C_j^{5_2}$) live on the $D5_2$-brane in the "bulk" between the source and matter branes. The gauge-singlet source field in principle can either live on the $D5'_1$-brane or be localised on the $5'_1 5_2$ intersection, but for definiteness we assume the latter possibility.

We now consider a limiting case in which the model in Fig. 3 reduces to the $\tilde{g}MSB$ model discussed in section 2, namely that the D5_2 radius of compactification is very much larger than the $D5_1$ radius[13]

$$R_{5_2} \gg R_{5_1} \gg m_s^{-1} \tag{8}$$

In this limit, the model reduces to that shown in Fig. 1, where the D3-branes correspond to the intersection regions of the D5-branes, and the bulk corresponds to the mediating 5_2 brane, as shown in Fig. 4. Note that the first two families are located on the matter brane, while the third family and Higgs doublets live on the mediating 5_2 brane.

[13]Both inverse radii must be larger than the inverse string scale.

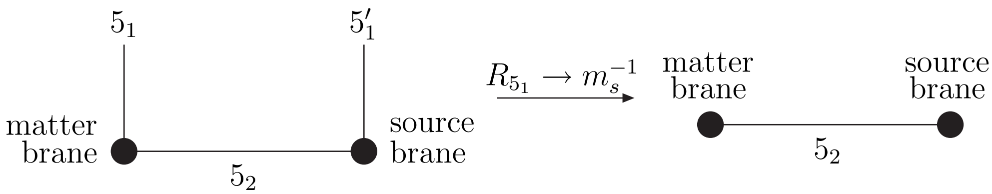

Figure 4: The intersecting D5-brane construction in the limit of small $D5_1$ compactification radius. The $D5_1$-branes reduce to effective D3-branes, separated in two orthogonal dimensions along the $D5_2$-brane (*bulk*). The allocation of Higgs and chiral matter fields are the same as in Fig. 3 and Table 2.

Since the gauge couplings on the branes are given by

$$g_{5_2}^{-2} = \frac{m_s^2 v_2}{(2\pi)^3 \lambda}, \qquad g_{5_1}^{-2} = \frac{m_s^2 v_1}{(2\pi)^3 \lambda} \tag{9}$$

we know that the coupling-squared is inversely proportional to compactification volume ($v_i \sim (2\pi R_i)^2$), which implies that $g_{5_1} \gg g_{5_2}$. This limiting case of the symmetry breaking is discussed in Appendix A.2, but the important results are that the dominant components of the gauge fields live on the $D5_2$-brane which is consistent with $\tilde{g}MSB$ with two extra bulk dimensions. After the gauge symmetry is broken down to the Standard Model, we recover the relationship between gauge couplings:

$$g_Y' \sim \sqrt{\frac{3}{5}} g_{5_2} \tag{10}$$

(where $g_3 \equiv g_{5_2}$). This is consistent with gauge coupling unification if $g_{5_2} \equiv g_{GUT}$ at the GUT scale.

It is also interesting to note that the restrictions we place on the radii do not restrict the radius of the third complexified dimension too strongly. This could allow a large extra dimension felt by gravity alone (with a size of the order of 1mm) as considered recently [14], but we will not discuss that possibility here.

In this limiting case, we can use the results of Ref. [4] where we identify $L \equiv R_{5_2}$. We can extend the analysis for the size of the extra dimensions and exponential suppression factors. Ref. [4] considers the maximum dimension size in the strong coupling limit $\epsilon \sim 1$, but for a small number of extra dimensions, the theory does not need to be strongly coupled at the string

scale, ie. $\epsilon \neq 1$.

Consider our symmetric toroidal compactification where the volume of the compact dimensions is

$$V_{5_2} \sim L^2 \equiv R_{5_2}^2 \tag{11}$$

Using eqs. (2,3) with $D = 6$, we can relate dimension size to ϵ for a 4-dimensional gauge coupling of order 1 (as observed for SM couplings).

$$\begin{aligned} g_{5_2}^2 &\sim \frac{\epsilon l_6}{m_s^2} \sim L^2 \\ L m_s &\sim (\epsilon l_6)^{\frac{1}{2}} \end{aligned} \tag{12}$$

Note that from eqs. (11, 12), we have

$$V_{5_2} m_s^2 \sim \epsilon l_6 \tag{13}$$

ϵ	Lm_s	$e^{-Lm_s/2}$
1	63	2×10^{-14}
0.8	56	6×10^{-13}
0.6	49	3×10^{-11}
0.4	40	2×10^{-9}
0.2	28	8×10^{-7}
0.1	20	5×10^{-5}
0.05	14	9×10^{-4}
0.01	6	4×10^{-2}

Table 3: Estimates for the toroidal compactification length L and exponential suppression factor for $D = 6$, where $L \equiv R_{5_2}$.

We have just seen how to recover the $\tilde{g}MSB$ model, but with two extra dimensions and the third family in the bulk. We can therefore use the $\tilde{g}MSB$ results for the operators that lead to scalar and Higgs masses, A and $B\mu$-terms and even a μ-term via the Giudice-Masiero mechanism[14]. However, in our model with $M \equiv m_s$ and $\mathrm{R}_{5_2} \gg R_{5_1} > m_s^{-1}$, there are only two extra dimensions in the bulk between D3-branes[15].

[14]Remember that a superpotential μ-term is forbidden by string selection rules for *our* choice of states.
[15]This allows us to use Table 3 to get restrictions on the size of R_{5_2}.

We use the eqs. (4,5,6) with the following identifications:

$$\begin{aligned}
\phi_M &\equiv \left(C^{5_1 5_2}\right) \quad Q_i, L_i, U_i^C, D_i^C, E_i^C, N_i^C \qquad (i=1,2) \\
W^\alpha &\equiv W_{SM} \\
h_u, h_d &\equiv \left(C_j^{5_2}\right) \quad h_u, h_d, Q_3, L_3, U_3^C, D_3^C, E_3^C, N_3^C \\
\hat{\phi}_S &\equiv \left(C^{5_1' 5_2}\right) \quad S
\end{aligned} \tag{14}$$

to generate higher-dimensional operators, subject to the full 42222 gauge invariance. We assume that the F-component of the gauge-singlet field S, which we assume to be an open string state on the intersection between the souce brane and the mediating brane, acquires a non-zero vev and breaks supersymmetry. We now proceed to discuss the different types of masses in the limiting case of the BMSB model.

4.1 Gaugino masses

In the limit of $R_{5_1} < R_{5_2}$, the Standard Model gauge fields are dominated by their components on the $D5_2$-brane (bulk). In agreement with $\tilde{g}MSB$, we generate gaugino masses of the same order of magnitude from eqs. (5,6)

$$m_\lambda \sim \frac{F_S}{m_s} \frac{l_6/l_4}{m_s^2 V_{5_2}} \sim \frac{1}{\epsilon l_4} \frac{F_S}{m_s} \tag{15}$$

(where V_{5_2} is the volume of the compact dimensions inside the $D5_2$-brane world-volume.)

4.2 First and second family scalar masses

This is the generic 4-point contact term between fields on opposite branes that leads to exponentially suppressed first and second family squark and slepton masses[16], using eq. (4).

$$\triangle \mathcal{L}_{soft} \sim \frac{e^{-m_s R_{5_2}}}{m_s^2} \int d^4\theta C^{\dagger 5_1 5_2} C^{5_1 5_2} S^\dagger S \tag{16}$$

$$V_{soft} \sim \frac{e^{-m_s R_{5_2}} F_S^2}{m_s^2} \left(\tilde{Q}_i^* \tilde{Q}_j + \tilde{U}_i^{C*} \tilde{U}_j^C + \tilde{D}_i^{C*} \tilde{D}_j^C + \tilde{L}_i^* \tilde{L}_j + \tilde{E}_i^{C*} \tilde{E}_j^C + \tilde{N}_i^{C*} \tilde{N}_j^C\right) \tag{17}$$

[16]This operator also leads to first and second family mixing and off-diagonal mass matrix elements. There may be another operator leading to first and third family mixing, eg. $\triangle \mathcal{L} \sim \int d^4\theta C^{\dagger 5_1 5_2} C_j^{5_2} S^\dagger S$.

Table (3) shows that the exponential suppression factor is strong for two extra dimensions. Therefore, contact term contributions to the first and second family scalar masses are negligible at high energies. Instead, they are generated by Renormalization Group Equation (RGE) effects.

Loop contributions to first and second family scalar masses (Fig. 1) are much larger than contact terms and anomaly mediated contributions. So, although the first and second family squark/slepton masses are not zero at high-energies, they are suppressed by a loop factor relative to third family scalar masses.

4.3 Higgs mass terms and third family scalar masses

Extending eq. (5) to include third family scalars, we have the following higher-dimensional operators:

$$\triangle\mathcal{L}_{soft} \sim \frac{l_6}{l_4}\int d^4\theta \left\{ \frac{1}{m_s^3}\left(S^\dagger h_u h_d + h.c.\right) + \frac{1}{m_s^4} S^\dagger S \left[h_u^\dagger h_u + h_d^\dagger h_d \right.\right.$$
$$\left.\left. + (h_u h_d + h.c.) + Q_3^\dagger Q_3 + U_3^{C\dagger} U_3^C + D_3^{C\dagger} D_3^C + L_3^\dagger L_3 + E_3^{C\dagger} E_3^C + N_3^{C\dagger} N_3^C \right] \right\} \tag{18}$$

From eqs. (6, 18), we obtain the μ-term,

$$\mu \sim \frac{F_S}{m_s}\frac{l_6/l_4}{m_s^2 V_{5_2}} \sim \frac{1}{\epsilon l_4}\frac{F_S}{m_s} \tag{19}$$

Higgs and third family scalar masses.

$$B\mu, m_{h_u}^2, m_{h_d}^2, m_{\tilde{F}_3}^2 \sim \frac{F_S^2}{m_s^2}\frac{l_6/l_4}{m_s^2 V_{5_2}} \sim \frac{1}{\epsilon l_4}\frac{F_S^2}{m_s^2} \tag{20}$$

(where $\tilde{F}_3 \equiv \tilde{Q}_3, \tilde{U}_3^C, \tilde{D}_3^C, \tilde{L}_3, \tilde{E}_3^C, \tilde{N}_3^C$)

4.4 Scalar mass matrix

We have generated a scalar mass matrix with an explicit third family mass hierarchy at lowest order:

$$m_{scalar}^2 \sim \frac{1}{\epsilon l_4}\frac{F_S^2}{m_s^2}\begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \tag{21}$$

The first and second family mass matrix elements are dominated by loop corrections since the contact term contributions are exponentially suppressed. However these contributions are still smaller than the third family masses due to the location of the third family in the bulk and its direct coupling to the SUSY breaking hidden sector.

4.5 Trilinear A-terms

Gauge invariant operators can be constructed for third family A-terms as follows:

$$\triangle\mathcal{L}_{soft} \sim \frac{l_6}{l_4}\int d^2\theta \frac{1}{m_s^4} S\left(h_d D_3^C Q_3 + h_u U_3^C Q_3 + h_d E_3^C L_3 + h_u N_3^C L_3\right) + h.c. \tag{22}$$

These operators lead to trilinear A-terms:

$$A_{ij} \sim \frac{F_S}{m_s}\frac{l_6/l_4}{m_s^3 V_{5_2}^{3/2}}\begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \sim \frac{F_S}{m_s}\frac{1}{\epsilon l_4\left(\epsilon l_6\right)^{1/2}}\begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \tag{23}$$

using eq. (13).

The first and second family A-terms are negligible in comparison to the third family term. Instead, they will receive loop-suppressed contributions.

4.6 Yukawa textures

Using our choice of states and eq. (7), we obtain a third family hierarchical Yukawa texture for the quark and lepton sectors at lowest-order. This texture reflects the observation that $\mathrm{m}_t \gg m_c, m_u$; $\mathrm{m}_b \gg m_s, m_d$ and $\mathrm{m}_\tau > m_\mu, m_e$.

$$Y_{ij}^a \sim \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \text{ where } a \equiv u, d, e, n \tag{24}$$

Smaller NLO Yukawa couplings (and associated trilinear A-terms) are generated by higher-dimensional operators. Notice that an interesting operator is allowed by 42222 gauge invariance, and appears to be such a small Yukawa term:

$$\triangle\mathcal{L} \sim \bar{F}_i F_i h \phi \phi' \tag{25}$$

The fields h, ϕ and ϕ' (Higgs) acquire vevs and spontaneously break the gauge symmetry. When each field is replaced by its vev, we can generate a first and second family mass term. This operator will be suppressed by powers of the string scale such that the first and second family have much smaller masses relative to the third family in the bulk.

4.7 Mass ratios and FCNC constraints

Consider the ratio of Higgs and third family scalar masses $B\mu, m_{h_u}^2, m_{h_d}^2, m_{\tilde{F}_3}^2$ to gaugino masses m_λ^2:

$$\frac{m_\phi^2}{m_\lambda^2} \sim \frac{l_4}{l_6} m_s^2 V_{5_2} \sim \epsilon l_4 \tag{26}$$

(using eqs. (2,3) and $g_4 \sim 1$, where $\mathrm{m}_\phi^2 \equiv B\mu, m_{h_u}^2, m_{h_d}^2, m_{\tilde{F}_3}^2$)

Also consider the ratio of trilinear soft masses A_{33} to gaugino masses m_λ using eqs. (15,23):

$$\frac{A_{33}}{m_\lambda} \sim \frac{1}{(\epsilon l_6)^{1/2}} \tag{27}$$

ϵ	$e^{-Lm_s/2}$	m_ϕ^2/m_λ^2	m_ϕ/m_λ	A_{33}/m_λ
1.0	2×10^{-14}	158	12.6	0.016
0.8	6×10^{-13}	126	11.2	0.018
0.6	3×10^{-11}	95	9.7	0.020
0.4	2×10^{-9}	63	7.9	0.025
0.2	8×10^{-7}	32	5.6	0.035
0.1	5×10^{-5}	16	4.0	0.050
0.05	9×10^{-4}	8	2.8	0.071
0.01	4×10^{-2}	1.6	1.3	0.159

Table 4: Estimates for the ratio of scalar masses and third family A-terms to gaugino masses for different ϵ and the exponential suppression factor (for masses-squared) arising from toroidal compactification.

Experimental constraints on FCNC[17] from mass-squared matrix elements require an exponential suppression of $\sim 10^{-3} - 10^{-4}$ for first and second family scalar masses in eq. (17). Using Table 4, we get a *lower* limit of say $\epsilon \sim 0.01$. However, phenomenological considerations restrict

[17]See [4] and references therein.

the ratio of m_ϕ and m_λ, and places an *upper* limit of say $\epsilon \sim 0.1$. This amount of suppression requires that the effective D3-branes are separated by a distance of order $\sim 10/m_s$.

4.8 Phenomenology

As in [4] we shall consider the phenomenology based on an inverse compactification scale ($R_{5_2}^{-1}$ in our case) close to the unification scale $M_{GUT} \sim 2 \times 10^{16}$ GeV. It is natural to assume a high energy unification scale in the limiting case $g_{5_1} \gg g_{5_2}$ since in this limit the light physical gauge fields all arise from the mediating 5_2 brane, and so are all subject to a single gauge coupling constant, $g_{5_2} \equiv g_{GUT}$.

We have seen that in the BMSB model (at M_{GUT}) the trilinear and first and second family soft masses are negligible, while the third family soft masses, and the Higgs mass parameters are larger than the gaugino masses. In Table 5 we compare a sample spectrum in the BMSB model to that in both the $\tilde{g}MSB$ model and the no-scale supergravity model, where the ratio of Higgs vevs $\tan\beta = 20$ and a universal gaugino mass of $M_{1/2} = 300$ GeV are chosen to give a lightest Higgs boson mass of about 115 GeV, consistent with the recent LEP signal [15, 16]. [18]

In the no-scale model the only non-zero soft mass is $M_{1/2}$, which results in a very characteristic spectrum where the right-handed slepton is very light and is in danger of becoming lighter than the lightest neutralino. The $\tilde{g}MSB$ model differs from the no-scale model only by the inclusion of Higgs soft masses which we have taken to be degenerate and somewhat higher than the gaugino masses. The main effect is to reduce the μ parameter, which is determined here from the electroweak symmetry breaking condition, and taken to be positive, which results in lighter charginos and neutralinos. Also in the $\tilde{g}MSB$ model the heavy Higgs and third family squark spectrum is also noticeably different from the no-scale model. [19] Turning to the BMSB model, we see that the effect of having both the Higgs and third family soft masses is to raise the μ parameter, and of course to significantly increase the third family squark and slepton masses,

[18] Note that $\tan\beta = 20$ is sufficiently small that we may neglect all Yukawa couplings except the top Yukawa coupling in the RGEs.

[19] As noted in [4], if we had taken non-degenerate Higgs soft masses then the lightest right-handed slepton mass could have been significantly increased relative to the no-scale model due to the hypercharge Fayet-Illiopoulos term.

	BMSB	$\tilde{g}MSB$	no-scale
$M_{1/2}$	300	300	300
A_0	0	0	0
$m_{\tilde{F}_{1,2}}$	0	0	0
$m_{\tilde{F}_3}$	500	0	0
m_{h_u}	500	500	0
m_{h_d}	500	500	0
$\tilde{g}$	830	830	830
$\tilde{\chi}_1^0$	124	119	124
$\tilde{\chi}_2^0$	239	200	237
$\tilde{\chi}_3^0$	506	258	472
$\tilde{\chi}_4^0$	517	314	485
$\tilde{\chi}_1^\pm$	238	195	237
$\tilde{\chi}_2^\pm$	518	314	486
$\tilde{E}_{L_{1,2}}$	220	220	220
$\tilde{E}_{L_3}$	546	220	220
$\tilde{E}_{R_{1,2}}$	124	124	124
$\tilde{E}_{R_3}$	515	124	124
$\tilde{N}_{L_{1,2}}$	205	205	205
$\tilde{N}_{L_3}$	540	205	205
$\tilde{U}_{L_{1,2}}$	740	740	740
$\tilde{U}_{L_3}$	783	653	676
$\tilde{U}_{R_{1,2}}$	715	715	715
$\tilde{U}_{R_3}$	628	520	577
$\tilde{D}_{L_{1,2}}$	744	744	744
$\tilde{D}_{L_3}$	787	658	681
$\tilde{D}_{R_{1,2}}$	713	713	713
$\tilde{D}_{R_3}$	871	713	713
$\tilde{t}_1$	613	492	544
$\tilde{t}_2$	832	718	745
$\tan\beta$	20	20	20
m_{h^0}	115	114	115
m_{H^0}	738	596	511
m_A	738	596	511
$m_{H^\pm}$	742	602	517
$\mu(M_Z)$	500	250	467

Table 5: Comparison of spectra (in GeV) for the three models BMSB, $\tilde{g}MSB$ and no-scale supergravity. The common parameters are $\tan\beta = 20$, universal gaugino mass $M_{1/2} = 300$ GeV, trilinear soft mass $A_0 = 0$, first and second family squark and slepton masses $m^2_{\tilde{F}_{1,2}} = 0$. The parameters are chosen to give a lightest Higgs boson mass consistent with the LEP signal [15, 16]. The μ parameter (assumed positive) and B are determined from the low energy electroweak symmetry breaking conditions.

providing an unmistakable spectrum and a characteristic smoking gun signature of the model.

5 Conclusions

We have proposed a mechanism for mediating SUSY breaking in Type I string theories - BMSB. Rather similar to the $\tilde{g}MSB$ set-up in Fig.1 we have proposed a Type I string-inspired set-up consisting of three intersecting D5-branes as shown in Fig.3 in which the gauge fields, Higgs doublets and third family matter fields all live on the third mediating 5_2-brane which plays the role of the bulk in the $\tilde{g}MSB$ scenario. The presence of the third matter family on the mediating D-brane is characteristic of Type I string constructions and provides the main experimentally testable difference between the BMSB and $\tilde{g}MSB$ models.

We have considered a limiting case in which $R_{5_2} \gg R_{5_1}$, and shown that in this case the model reduces to the original $\tilde{g}MSB$ model with the role of the bulk being played by the mediating 5_2-brane. In this limiting case, the model naturally leads to approximately universal gaugino masses and a single unified gauge coupling constant, which motivates the identification of the string scale with the usual GUT scale. In this case the phenomenology of the BMSB model is rather interesting, and it may be compared to the predictions of the no-scale supergravity and the $\tilde{g}MSB$ model. As in the $\tilde{g}MSB$ model, the first two families naturally receive very small masses at the high energy scale leading to flavour-changing neutral currents being naturally suppressed. The presence of third family soft masses will not alter this conclusion very much since FCNC limits involving the third family are much weaker. However the third family soft masses will lead to a characteristic squark and slepton mass spectrum which may be easily distinguished from that of both no-scale supergravity and the $\tilde{g}MSB$ model as shown in Table 5. The μ-problem is solved by the Giudice-Masiero mechanism as in the original $\tilde{g}MSB$ model.

In this limiting case the BMSB model bears a close resemblance to both the no-scale supergravity and the $\tilde{g}MSB$ models. The fact that the third family receives a non-zero soft SUSY breaking mass is strictly not an unambiguous signal of the underlying Type I string model, since it is possible for this to happen in both the other cases also. For example in the old heterotic

models based on orbifolds, matter may be localised in the fixed points of the orbifold (the twisted sector) or not (the untwisted sector) so it is possible to have the third family playing a different role from that of the first two families. What is different in the model presented here is that the gauge group is localised on two different branes, but in the limiting case (above) the physical gauge group arises essentially from one brane, and in this limit we return to a situation similar to that of the old heterotic string theories. There are however three points worth noting here. Firstly, the presence of two families at the intersection points of two branes, and one family on a single brane, seems to be typical of Type I string constructions [8]. Secondly in Type I string constructions we have the possiblity of full unification of both gravity and gauge forces, precisely because gravity exists in 10 dimensions whereas the gauge groups live in a 6 dimensional sub-manifold, which is not possible in old heterotic string theories. Thirdly, the limiting case of $R_{5_2} \gg R_{5_1}$ would be expected to apply only approximately, and therefore in practice there will be corrections, for example, to gauge coupling unification which may be observable. Further comments concerning the non-limiting case are briefly discussed below.

In the more general non-limiting case, the model will have an even richer structure. In this non-extremal radii limit (ie. $\mathrm{R}_{5_2} > R_{5_1} \neq m_s^{-1}$), we must use the full gauge state expressions listed in Appendix A.1. The light gauge states are no longer dominated by their $D5_2$-brane components, but are instead mixtures of fields from either brane, with the exception of the gluon/gluino states that *only* arise from the $D5_2$-brane. The result is that the high energy gluino mass will be larger than the high energy wino and bino masses. In this more general case the gauge couplings are no longer equal, so there is less motivation to identify the string scale with the GUT scale. Generally the string scale can take any value from a few TeV to 10^{16} GeV, and we have the possibility of a mm scale large extra dimension.

The toy model has other interesting features such as the fact that the gauge symmetry forbids first and second family Yukawa couplings at lowest order, and naturally forbids R-parity violating operators that cannot be forbidden by string selection rules alone, while allowing the third family Yukawa coupling. Most importantly, however, the toy model demonstrates the

BMSB mechanism, which is based on having at least three branes with two different intersection points. This minimum requirement implies that constructions with all the branes at the origin fixed point are inadequate for our purpose. Although there are examples in the literature of intersecting branes at different fixed points [17], such models are generally more complicated than the simple set-up considered here. Nevertheless our BMSB mechanism could provide a useful alternative starting point from which to address the problem of SUSY breaking in more general Type I string theories.

Finally note that it has been been suggested, in the context of Type I theories, that singlet twisted moduli, which appear in the tree-level gauge kinetic function, might be responsible for generating gaugino masses if they acquire non-vanishing F-terms, and that this might provide a brane realisation of $\tilde{g}MSB$ if the standard model gauge symmetry originates from 9-branes providing that there are in addition two sets of D5-branes located at two different fixed points [18]. This suggestion shares some of the features with the present paper, although model building issues were not discussed, and the characteristic possibility of the third family on the mediating brane was not considered. Also additional contributions from the F-terms of dilaton S and moduli fields T_i were also generically allowed, whereas here we have implicitly assumed them to be absent.

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A Spectrum of Gauge bosons

A.1 General case

In this appendix we consider the effect of symmetry breaking on massless gauge field states and gauge couplings. We begin with the gauge group $SU(4)_{5_2} \otimes SU(2)_{5_{1L}} \otimes SU(2)_{5_{2L}} \otimes SU(2)_{5_{1R}} \otimes SU(2)_{5_{2R}}$. The couplings run with energy scales subject to RGEs. Conventionally, the symmetry breaking all occurs at high energies $(10^{15} - 10^{16} GeV)$ except for $\mathrm{SU}(2)_L \otimes U(1)_Y \longrightarrow U(1)_{EM}$ which happens at the Electroweak scale. In the tables that follow, gauge couplings are assumed to be at high energies unless otherwise stated. Notice that i, a and m are adjoint indices for SU(2), SU(3) and SU(4) respectively.

Gauge group	$SU(4)_{5_2}$	$SU(2)_{5_{1L}}$	$SU(2)_{5_{2L}}$	$SU(2)_{5_{1R}}$	$SU(2)_{5_{2R}}$
Coupling	g_{5_2}	g_{5_1}	g_{5_2}	g_{5_1}	g_{5_2}
States	$G^m_{5_2}$	$W^i_{5_1 L}$	$W^i_{5_2 L}$	$W^i_{5_1 R}$	$W^i_{5_2 R}$

Table 6: The initial gauge groups, gauge couplings and states in our model.

(a) First combine the chiral SU(2) groups from either brane via diagonal symmetry breaking to recover the Pati-Salam gauge group.

$$SU(2)_{5_{1L/R}} \otimes SU(2)_{5_{2L/R}} \overset{diagonal}{\xrightarrow{v_\phi, v_{\phi'}}} SU(2)_{L/R}$$
$$\Rightarrow SU(4)_{5_2} \otimes SU(2)_L \otimes SU(2)_R \equiv G_{PS} \tag{28}$$

Spontaneous symmetry breaking (SSB) induces a change of basis, parametrised by

$$\cos\theta_\phi = \frac{g_{5_2}}{\sqrt{g_{5_1}^2 + g_{5_2}^2}} \tag{29}$$

We can express the new massless states and gauge couplings in terms of the original parameters. The Higgs mechanism generates massive gauge bosons with masses of the order of the symmetry breaking scale.

Gauge group	$SU(4)_{5_2}$	$SU(2)_L$ $\quad$ $SU(2)_R$
Coupling	g_{5_2}	$g_L = \frac{g_{5_1} g_{5_2}}{\sqrt{g_{5_1}^2 + g_{5_2}^2}} = g_R$
States	$G_{5_2}^m$	$W_{L/R}^i = \frac{1}{\sqrt{g_{5_1}^2 + g_{5_2}^2}} \left(g_{5_1} W_{5_2 L/R}^i + g_{5_2} W_{5_1 L/R}^i \right)$

Table 7: The new massless states and couplings after the original gauge symmetry is broken down to the Pati-Salam gauge group.

plus 3 massive $SU(2)_L$ $\left(\bar{W}_L\right)$ and 3 massive $SU(2)_R$ $\left(\bar{W}_R\right)$ bosons, of mass

$$M_{\bar{W}_{L/R}}^2 = \tfrac{1}{2} v_\phi^2 \left(g_{5_1}^2 + g_{5_2}^2 \right)$$

(b) QCD $SU(3)_C$ is contained within $SU(4)_{5_2}$. The U(1)s combine to give the Hypercharge U(1) using the relationship $Y = (B-L) + 2I_R$.

$$\begin{aligned} SU(4)_{5_2} &\supset SU(3)_C \otimes U(1)_{B-L} \\ SU(2)_R &\supset U(1)_{I_R} \end{aligned} \tag{30}$$

The Pati-Salam gauge group is broken down to the Standard Model by giving a vev to a Higgs field H.

$$\begin{aligned} & U(1)_{B-L} \otimes U(1)_{I_R} \xrightarrow{v_H} U(1)_Y \\ & \Rightarrow SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \end{aligned} \tag{31}$$

The change of basis is parametrised by

$$\cos\theta_H = \frac{\sqrt{\frac{3}{2}} g_{5_2}}{\sqrt{g_R^2 + \frac{3}{2} g_{5_2}^2}} = \sqrt{\frac{3\left(g_{5_1}^2 + g_{5_2}^2\right)}{5g_{5_1}^2 + 3g_{5_2}^2}} \tag{32}$$

Gauge group	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
Coupling	g_{5_2}	$g_L = \dfrac{g_{5_1} g_{5_2}}{\sqrt{g_{5_1}^2 + g_{5_2}^2}}$	$g_Y^{'} = \dfrac{g_{5_1} g_{5_2} \sqrt{3}}{\sqrt{5g_{5_1}^2 + 3g_{5_2}^2}}$
States	$G_{5_2}^a$	$W_L^i = \dfrac{g_{5_1} W_{5_2 L}^i + g_{5_2} W_{5_1 L}^i}{\sqrt{g_{5_1}^2 + g_{5_2}^2}}$	$B_Y = \dfrac{\sqrt{3}\left(g_{5_1} W_{5_2 R}^3 + g_{5_2} W_{5_1 R}^3\right) + \sqrt{2} g_{5_1} G_{5_2}^{15}}{\sqrt{5g_{5_1}^2 + 3g_{5_2}^2}}$

Table 8: The Standard Model massless states and gauge couplings expressed in terms of the original parameters.

$$\text{plus 6 massive } SU(4)_{5_2} \text{ bosons } \left(G_{5_2}^9 - G_{5_2}^{14}\right), \text{ mass } M_G^2 = \tfrac{1}{4} v_H^2 g_{5_2}^2,$$
$$2 \text{ massive } SU(2)_R \text{ bosons } \left(W_R^{\pm}\right), \text{ mass } M_{W_R^{\pm}}^2 = \tfrac{1}{4} v_H^2 g_R^2,$$
$$\text{and 1 massive } SU(2)_{B-L} \text{ boson } (X_{B-L}), \text{ mass } M_{X_{B-L}}^2 = \tfrac{1}{4} v_H^2 \left(g_R^2 + \tfrac{3}{2} g_{5_2}^2\right)$$

(c) Finally, we can recover the QCD and EM Standard Model gauge group via the familiar low-energy Higgs mechanism, parametrised by

$$\cos\theta_W = \frac{g_L(v_h)}{\sqrt{g_L^2(v_h) + g_Y^{'2}(v_h)}} = \sqrt{\frac{5g_{5_1}^2(v_h) + 3g_{5_2}^2(v_h)}{8g_{5_1}^2(v_h) + 6g_{5_2}^2(v_h)}}.$$

Electroweak symmetry breaking occurs when the Higgs field h acquires a non-zero vev.

$$\begin{aligned} SU(2)_L \otimes U(1)_Y &\xrightarrow{v_h} U(1)_{EM} \\ &\Rightarrow SU(3)_C \otimes U(1)_{EM} \end{aligned} \tag{33}$$

plus 3 massive $SU(2)_L$ bosons $(W_L^{\pm}, Z_L^0)$ with masses:

$$M_{W_L^{\pm}} = \frac{1}{2} v_h g_L(v_h) = \frac{g_{5_1}(v_h) g_{5_2}(v_h) v_h}{2\sqrt{g_{5_1}^2(v_h) + g_{5_2}^2(v_h)}} \text{ and}$$

$$M_{Z_L^0} = g_{5_1}(v_h) g_{5_2}(v_h) v_h \sqrt{\frac{4g_{5_1}^2(v_h) + 3g_{5_2}^2(v_h)}{2\left(g_{5_1}^2(v_h) + g_{5_2}^2(v_h)\right)\left(5g_{5_1}^2(v_h) + 3g_{5_2}^2(v_h)\right)}}$$

Gauge group	$SU(3)_C$	$U(1)_{EM}$
Coupling	$g_{5_2}(v_h)$	$e = \frac{g_{5_1}(v_h) g_{5_2}(v_h)\sqrt{3}}{\sqrt{8g_{5_1}^2(v_h) + 6g_{5_2}^2(v_h)}}$
States	$G_{5_2}^a$	$A = \frac{\sqrt{3}g_{5_1}(v_h)\left(W_{5_2L}^3 + W_{5_2R}^3\right) + \sqrt{3}g_{5_2}(v_h)\left(W_{5_1L}^3 + W_{5_1R}^3\right) + \sqrt{2}g_{5_1}(v_h)G_{5_2}^{15}}{\sqrt{8g_{5_1}^2(v_h) + 6g_{5_2}^2(v_h)}}$

Table 9: The massless gauge states and couplings after electroweak symmetry breaking.

A.2 Limiting case $R_{5_2} \gg R_{5_1}$

In this appendix we repeat the symmetry breaking analysis for the limiting case

$$R_{5_2} \gg R_{5_1} \Leftrightarrow g_{5_2} \ll g_{5_1} \tag{34}$$

We find that the dominant components of the massless gauge fields live on the 5_2-brane ("bulk") which is consistent with $\tilde{g}MSB$.

(a) After diagonal symmetry breaking we recover the Pati-Salam gauge group

S$U(4)_{5_2} \otimes SU(2)_L \otimes SU(2)_R$

Gauge group	$SU(4)_{5_2}$	$SU(2)_L$ $SU(2)_R$
Coupling	g_{5_2}	$g_{L/R} \sim g_{5_2}$
States	$G_{5_2}^m$	$W_{L/R}^i \sim W_{5_2L/R}^i$

Table 10: The dominant components of massless states and couplings after symmetry has been broken down to the Pati-Salam group.

$$\text{plus 3 massive } SU(2)_L \text{ and 3 massive } SU(2)_R \text{ bosons } \left(\bar{W}_{L/R} \sim W_{5_1L/R}^i\right),$$
$$M_{\bar{W}_{L/R}}^2 \approx \tfrac{1}{2} v_\phi^2 g_{5_1}^2$$

(b) We break the Pati-Salam group down to the Standard Model. Notice the relationship between the Hypercharge gauge coupling and the other gauge couplings, which is consistent with gauge coupling unification. This will happen if the 5_2 gauge coupling equals g_{GUT} at the GUT scale.

Gauge group	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
Coupling	g_{5_2}	$g_L \sim g_{5_2}$	$g_Y^{'} \sim \sqrt{\frac{3}{5}} g_{5_2}$
States	$G^a_{5_2}$	$W^i_L \sim W^i_{5_2 L}$	$B_Y \sim \sqrt{\frac{3}{5}} W^3_{5_2 R} + \sqrt{\frac{2}{5}} G^{15}_{5_2}$

Table 11: The dominant components of the massless states and couplings after the Pati-Salam group is broken down to the Standard Model.

$$\text{plus 6 massive } SU(4)_{5_2} \text{ bosons } \left(G^9_{5_2} - G^{14}_{5_2}\right),\ M^2_G \approx \tfrac{1}{4} v^2_H g^2_{5_2},$$
$$\text{2 massive } SU(2)_R \text{ bosons } \left(W^\pm_R \sim \tfrac{1}{\sqrt{2}}\left(W^1_{5_2 R} \mp i W^2_{5_2 R}\right)\right),\ M^2_{W^\pm_R} \approx \tfrac{1}{4} v^2_H g^2_{5_2}$$
$$\text{and 1 massive } SU(2)_{B-L} \text{ boson } \left(X_{B-L} \sim \sqrt{\tfrac{3}{5}} G^{15}_{5_2} - \sqrt{\tfrac{2}{5}} W^3_{5_{2R}}\right),\ M^2_{X_{B-L}} \approx \tfrac{5}{8} g^2_{5_2} v^2_H$$

(c) Finally the Higgs mechanism induces electroweak symmetry breaking, and generates the massive W and Z bosons.

Gauge group	$SU(3)_C$	$U(1)_{EM}$
Coupling	$g_{5_2}(v_h)$	$e \sim \sqrt{\frac{3}{8}} g_{5_2}(v_h)$
States	$G^a_{5_2}$	$A \sim \sqrt{\frac{3}{8}} \left(W^3_{5_2 L} + W^3_{5_2 R}\right) + \frac{1}{2} G^{15}_{5_2}$

Table 12: The dominant components of the familiar massless gauge states after electroweak symmetry.

$$\text{plus 3 massive } SU(2)_L \text{ bosons:}$$
$$\left(W^\pm_L \sim \tfrac{1}{\sqrt{2}}\left(W^1_{5_2 L} \mp i W^2_{5_2 L}\right)\right),\ M_{W^\pm_L} \approx \tfrac{1}{2} v_h g_{5_2}(v_h)$$
$$\text{and } \left(Z^0_L \sim \sqrt{\tfrac{5}{8}} W^3_{5_2 L} - \tfrac{3}{2\sqrt{10}} W^3_{5_2 R} - \tfrac{1}{2}\sqrt{\tfrac{3}{5}} G^{15}_{5_2}\right),\ M_{Z^0_L} \approx \sqrt{\tfrac{2}{5}} v_h g_{5_2}(v_h)$$

B Mass scales

In this appendix we consider the different mass scales present in the model. Each time the gauge symmetry is spontaneously broken down towards the Standard Model, the broken generators have massive gauge bosons associated with them. These bosons have masses of the same order as the symmetry breaking scale, ie. the vevs of the breaking fields. Our model already assumes

an order for symmetry breaking, which creates a vev hierarchy ($v_\phi \geq v_H \gg v_h \sim O\left(M_{EW}\right)$). For instance, we know that $v_\phi, v_H \gg O\left(M_{EW}\right)$ since these broken symmetry bosons have not been observed.

We must also consider the (inverse) compactification radii of the D5-branes. Their relative sizes are arbitrary, but we choose to start with the relationship $R_{5_2} > R_{5_1}$ or equivalently $R_{5_1}^{-1} > R_{5_2}^{-1}$ as shown in Fig. 5.

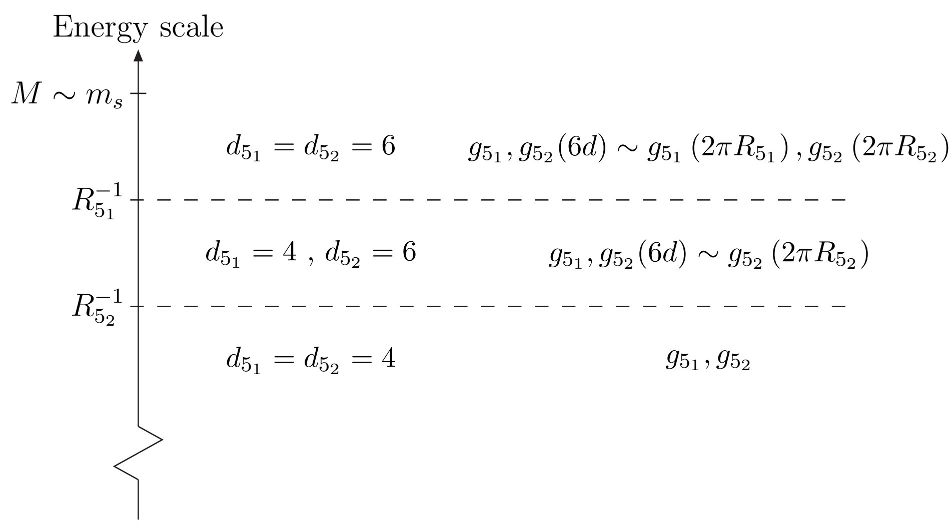

Figure 5: At energy scales below an inverse compactification radii, the dimension appears too small to observe. The coupling in a higher-dimension is related to the same coupling in a lower dimension via eq. (2).

Notice that we have not specified how R_{5_3} is related to the other two compactification radii, suffice to say that a large third dimension (felt by gravity alone) is not forbidden ie.$R_{5_1}^{-1} > R_{5_2}^{-1} \gg R_{5_3}^{-1}$ (See [14] for discussion of large extra dimensions).

In this work, we have adopted the standard scenario with symmetry breaking occurring at a scale comparable to the first two compactification radii and string scale. Soft masses are also generated at around the same scale. We have deliberately not specified these scales, but we claim that the formalism applies for GUT/string scales of $1TeV$ to $10^{16}GeV$.

We impose the following restrictions:

$$R_{5_1}^{-1} \quad > R_{5_2}^{-1}$$

$$\begin{aligned} v_\phi &\geq v_H \gg v_h \\ m_s &\geq R_{5_1}^{-1}, R_{5_2}^{-1} \sim v_\phi, v_H \gg v_h \sim O\left(M_{EW}\right) \end{aligned} \tag{35}$$

These constraints provide six ways of ordering the inverse radii and vevs. The supersymmetry breaking scale (where soft masses are generated) also needs to be assigned, thus giving a total of 30 possibilities.

A	v_h	v_H	v_ϕ	$R_{5_2}^{-1}$	$R_{5_1}^{-1}$	m_s
B	v_h	v_H	$R_{5_2}^{-1}$	v_ϕ	$R_{5_1}^{-1}$	m_s
C	v_h	$R_{5_2}^{-1}$	v_H	v_ϕ	$R_{5_1}^{-1}$	m_s
D	v_h	v_H	$R_{5_2}^{-1}$	$R_{5_1}^{-1}$	v_ϕ	m_s
E	v_h	$R_{5_2}^{-1}$	v_H	$R_{5_1}^{-1}$	v_ϕ	m_s
F	v_h	$R_{5_2}^{-1}$	$R_{5_1}^{-1}$	v_H	v_ϕ	m_s

Table 13: Possible ordering of symmetry breaking vevs and inverse compactification radii within the constraints of eq. (35).

In Table 13 we list the various possibilities for the relative ordering of mass scales, vevs and inverse compactification radii within the constraints of eq.(35). However it is important to notice that when the inverse compactification radii are less than the symmetry breaking vevs, the Kaluza-Klein modes associated with the extra dimensions can contribute to the running of gauge couplings, leading to a power law dependence [9].

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